Realistic two-nucleon potentials for the relativistic two-nucleon Schroedinger equation

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The potentials $V(v)$ in the nonrelativistic (relativistic) nucleon-nucleon (NN) Schrödinger equation are related by a quadratic equation. That equation is numerically solved, thus providing phase equivalent v - potentials related for instance to the high precision NN potentials, which are adjusted to NN phase shift and mixing parameters in a nonrelativistic Schrödinger equation. The relativistic NN potentials embedded in a three-nucleon (3N) system for total NN momenta different from zero are also constructed in a numerically precise manner. They enter into the relativistic interacting 3N mass operator, which is needed for relativistic 3N calculations for bound and scattering states.

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I. INTRODUCTION

Traditionally the (semi) phenomenological high precision two-nucleon (NN)potentials AV18 [1], CD Bonn [2] and Nijm I,II [3] go together with the nonrelativistic operator for the kinetic energy $\frac{\hat{k}^2}{m}$ in the NN c.m.system. Nevertheless, as is well known, this nonrelativistic Schrödinger equation

$$(\frac{\hat{k}^2}{m} + V)\Psi = \frac{k_0^2}{m}\Psi \tag{1}$$

can be related to an underlying relativistic NN Schrödinger equation

$$(2\sqrt{m^2+\hat{k}^2} + v)\Phi = 2\sqrt{m^2+k_0^2}\Phi \tag{2}$$

by a simple algebraic step[4],[5]. Applying $(2\sqrt{m^2+\hat{k}^2}+v)$ to (2) from the left one obtains

$$(4(m^2+\hat{k}^2) + 2\omega(\hat{k})v + 2v\omega(\hat{k}) + v^2)\Phi = 4(m^2+k_0^2)\Phi \tag{3}$$

which can be identically rewritten into (1) if one defines

$$V = \frac{1}{4m}(2\omega(\hat{k})v + 2v\omega(\hat{k}) + v^2) \tag{4}$$

with $\omega(\hat{k}) = \sqrt{m^2+\hat{k}^2}$. (We use $\hat{k}$ in order to distinguish the momentum operator from the number k). Therefore adjusting V in (1) to the NN phase shift and mixing parameters from a phase shift analysis and relating the c.m. momentum k_0 to the Lorentz invariant lab energy T_{lab} via

$$k_0 = \sqrt{\frac{mT_{lab}}{2}} \tag{5}$$

(a relation identically valid for relativistic and nonrelativistic kinematics) one has in fact solved a relativistic equation. We also see that Ψ equals Φ .The question remains, what is v given V? The formal solution of that quadratic equation (4) is

$$v = \sqrt{4mV + 4\omega(\hat{k})^2} - 2\omega(\hat{k}) \tag{6}$$

Why is v of interest? If one turns to the 3N system and would like to investigate relativistic effects [6],[7] the knowledge of v is very useful. It defines together with the relativistic kinetic energy the interacting NN mass operator, which is a key ingredient for building the interactive 3N mass operator [8]. Therefore we focus in this paper on the determination of v related to the high precision NN potentials via (4) or (6).

In [9],[10] a potential v has been determined directly fitting (2) to NN phase shifts. Thereby the Urbana v_{14} potential has been readjusted achieving a fair fit(though not of the quality of the high precision potentials). In [11] a momentum scale transformation

$$2m + \frac{k^2}{m} = 2\sqrt{m^2+q^2} \tag{7}$$

has been introduced which provides an analytical relation between V and v and guarantees that the S-matrix related to (1) at c.m.momentum k equals the S-matrix related to (2) at c.m. momentum q. In other words the relativistic and nonrelativistic S-matrices agree at the same energy. This,however, is misleading since the equality of the two S-matrices should hold at the same c.m.momenta [12]. A better, though still approximate approach to relate V and v has been given in [6].

On the other hand there is the possibility to add an interaction to the square of the free NN mass operator $h^2 \equiv (2\omega(\hat{k}))^2 + 4mV'$. Then $H \equiv \frac{h^2}{4m} - m$ has exactly the same form as (1) with $V' = V$ provided we identify h^2 with the mass operator of the interacting NN system [4], [6]. This is of course also obvious from the relations (1) - (4). The construction of the relativistic 3N Hamiltonian

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requires, however, the 3N mass operator h rather than its square. Therefore our aim here is to solve (4) and (6) exactly for v. This is outlaid in section II. The validity of the resulting v is verified by demonstrating that it provides exactly the same phase shift parameters using (2) as the underlying V using (1).

Next we regard the 3N mass operator where NN c.m. forces enter in the form [8]

$$v_p \equiv \sqrt{(2\omega(\hat{k})+v)^2+p^2} - \sqrt{(2\omega(\hat{k}))^2+p^2} \qquad (8)$$

with p the total NN momentum. The p-dependence arises since in a 3-body system the NN subsystems are not at rest. In section III we propose a simple manner to determine v_p in a numerical precise way. This opens now the door to use $v's$ which are equivalent to the underlying high precision potentials in a relativistic context in 3N bound and scattering problems. We summarize in section IV. A technical derivation is given in the Appendix.

II. THE POTENTIAL V

The determination of v using Eq. (6) can be achieved by a spectral decomposition. One can proceed in close analogy to the representation derived in [13] for v_p given in Eq. (8). We regard a specific partial wave state (or coupled ones) with given orbital angular momentum(a) , total spin and total angular momentum. For the sake of a simpler notation we will not show these quantum numbers explicitely. Using the completeness relation of bound and scattering states for the potential V one obtains

$$\begin{aligned} &< k|v|k\,'>=< k|\Psi_b > M_b < \Psi_b|k\,'> \\ &-2\omega(k)\frac{\delta(k-k')}{kk\,'} \\ &+\int_0^\infty dk''(k\,'')^2 < k|\Psi''_k > 2\sqrt{k''^2+m^2} < \Psi''_k|k\,'> \end{aligned} \qquad (9)$$

where $\Psi_b(k)$ is the nonrelativistic deuteron wave function of (1) and M_b the mass of the deuteron. Here we introduced a different definition of the binding energy, namely an implicit one:

$$M_b \equiv \sqrt{4m^2+4\epsilon_b\,'m} = 2m+\epsilon_b\,' - \frac{{\epsilon_b\,'}^2}{4m}+\cdots \qquad (10)$$

In lowest order it agrees with the usual one $M_b \equiv 2m+\epsilon_b$. This new definition of the binding energy has in addition the feature that it can naturally be written as $\epsilon_b\,' = -\frac{\kappa^2}{m}$ in agreement with the form of the energy eigenvalue of (1) at the bound state pole $k = i\kappa$ of the S-matrix.

The expression (9) can be identically rewritten into the form

$$\langle k|v|k'\rangle = \Psi_b(k) M_b \Psi_b(k')$$

$$\begin{aligned} &+ \frac{m}{k^2-k'^2}\{2\omega(k)\Re[T(k',k;\frac{k^2}{m})] \\ &- 2\omega(k\,')'\Re[T(k,k';\frac{k'^2}{m})]\} \\ &+ \frac{m^2}{k^2-k'^2}\times \\ &\{\ \mathcal{P}\int_0^\infty dk''(k\,'')^2\frac{2\omega(k\,'')}{k''^2-k^2} \\ &\quad T(k,k'';\frac{k''^2}{m})T^*(k',k'';\frac{k''^2}{m}) \\ &- \mathcal{P}\int_0^\infty dk''(k\,'')^2\frac{2\omega(k\,'')}{k''^2-k'^2} \\ &\quad T(k,k'';\frac{k''^2}{m})T^*(k',k'';\frac{k''^2}{m})\}. \end{aligned} \qquad (11)$$

where T is the standard NN T-matrix related to V via the nonrelativistic Lippmann Schwinger equation. The derivation of that form is defered to the Appendix. A numerical implementation has not yet been performed, but we expect no problem.

A second more simple way is to directly solve the quadratic operator equation (4). In momentum space it reads

$$\begin{aligned} &2m < k|V|k\,'> = (\omega(k)+\omega(k\,'))< k|v|k\,'> \\ &+\ \frac{1}{2}\int_0^\infty dk\,''k\,''^2 < k|v|k\,''>< k\,''|v|k\,'> \end{aligned} \qquad (12)$$

or

$$\begin{aligned} &< k|v|k\,'> \\ &+\ \frac{1}{2(\omega(k)+\omega(k\,'))}\int_0^\infty dk\,''k\,''^2 < k|v|k\,''>< k\,''|v|k\,'> \\ &= \frac{2m < k|V|k\,'>}{\omega(k)+\omega(k\,')} \end{aligned} \qquad (13)$$

We verified numerically that for all the realistic high precision potentials AV18, CD Bonn, and Nijm I,II the following very simple iterative scheme works

$$< k|v|k\,'>^{(0)} = \frac{2m < k|V|k\,'>}{\omega(k)+\omega(k\,')}, \qquad (14)$$

$$\begin{aligned} &< k|v|k\,'>^{(n+1)} = \frac{1}{2(\omega(k)+\omega(k\,'))}\times\{4m < k|V|k\,'> \\ &-\int_0^\infty dk\,''k\,''^2 < k|v|k\,''>^{(n)}< k\,''|v|k\,'>^{(n)}\} \end{aligned} \qquad (15)$$

with $n = 0, 1, \cdots$. For certain partial waves the iteration to converge requires an additional step, namely

$$v^{(n+1)} \xleftarrow{\text{redefine}} (av^{(n+1)}+bv^{(n)})/(a+b) \qquad (16)$$

where the constants a and b are typically 1.

We display in Table I an example documenting the convergence. In Figs 1-3 we show for an example the original nonrelativistic potential $V(k,k\,')$, the resulting relativistic potential $v(k,k\,')$ and the difference

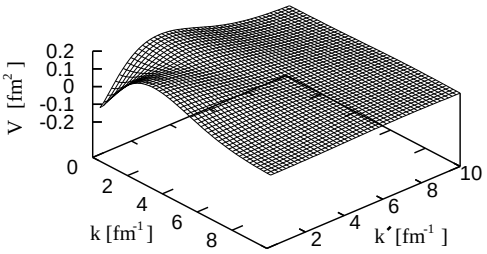

FIG. 1: The nonrelativistic potential $V(k,k')$ AV18 in the state $^1S_0$.

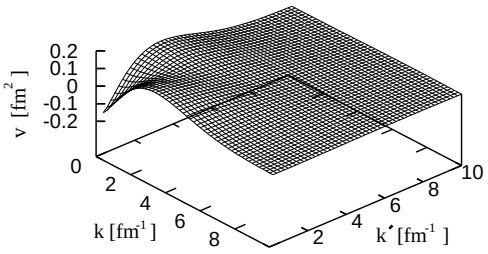

FIG. 2: The relativistic potential $v(k,k')$ related to AV18 in the state $^1S_0$.

$V(k,k\ ') - v(k,k\ ')$. We see in that example that $V(k,k\ ') > v(k,k\ ')$.

To further characterize the difference between $V(k,k\ ')$ and $v(k,k\ ')$ one can regard the asymptotic behavior

$$\lim_{k\to\infty, k\ '\text{fixed}} v(k,k\ ') = \frac{\text{const}}{k^n} \tag{17}$$

against

$$\lim_{k\to\infty, k\ '\text{fixed}} V(k,k\ ') = \frac{\text{const}}{k^N} \tag{18}$$

As examples , for AV18 we find $n = 3$, $N = 2$, whereas CD Bonn delivers $n = 4.5$, $N = 3.5$. The value of N for CD Bonn can easily be understood [14]: $N = 2+2+1/2-1$ with the two 2's resulting from the meson propagator and the choice of the strong form factor, the $1/2$ arising from transforming the Blankenbeclar-Sugar equation to the nonrelativistic Lippmann Schwinger equation and the (-1) from the two Dirac spinors. In both cases, AV18 and CD Bonn, n is larger than N by one unit, which is suggested by Eq.(12).

Having v at our disposal one can solve the relativistic Lippmann Schwinger equation

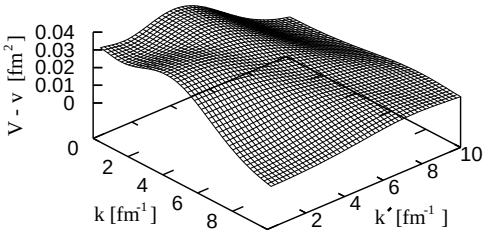

FIG. 3: The difference between the relativistic and nonrelativistic potentials: $V(k,k') - v(k,k')$ in the state $^1S_0$.

TABLE I: Convergence of $v^{(n)}$ to the iteration in Eq. (15). We choose the coupled partial waves ($^3S_1$-$^3D_1$) of the Argonne V18 potential[1]. The momenta k and k' are 1.0 fm^{-1} and the potential unit is [fm^2].

n	$v^{(n)}(^3S_1$-$^3S_1)$	$v^{(n)}(^3S_1$-$^3D_1)$	$v^{(n)}(^3D_1$-$^3D_1)$
0	0.084232	0.044709	0.016853
1	0.067716	0.044628	0.016785
2	0.059933	0.044597	0.016744
3	0.056135	0.044587	0.016719
4	0.054234	0.044585	0.016705
5	0.053259	0.044587	0.016696
6	0.052749	0.044589	0.016691
10	0.052194	0.044595	0.016684
20	0.052126	0.044597	0.016684
30	0.052126	0.044597	0.016684

$$\begin{aligned} t(k,k\ ') &= v(k,k\ ') \\ &+ \int_0^\infty dk\ ''k\ ''^2 v(k,k\ '') \\ &\frac{1}{2\omega(k\ ') - 2\omega(k\ '') + i\epsilon} t(k\ '',k\ ') \end{aligned} \tag{19}$$

for the half shell t-matrix. It is related to the S-matrix via

$$s(k) = e^{2i\delta_r(k)} = 1 - i\pi k\omega(k) t(k,k). \tag{20}$$

The corresponding nonrelativistic relation is

$$S(k) = e^{2i\delta_{nr}(k)} = 1 - i\pi k m T(k,k). \tag{21}$$

where $T(k,k\ ')$ obeys the standard nonrelativistic Lippmann Schwinger equation. We illustrate in Table II for the ($^3S_1$ - $^3D_1$) partial wave states the perfect agreement of the corresponding phase shift and mixing parameters.

TABLE II: Comparison of the phase shift and the mixing parameters
for the coupled partial waves ($^3S_1$-$^3D_1$) of the Argonne V18 potential[1]. The second column points to the relativistic (Rel.) or nonrelativistic (Nonrel.) calculations. The unit of the phases are in degrees.

$T_{lab.}$ [MeV]	Nonrel./Rel.	$\delta(^3S_1)$	$\delta(^3D_1)$	ϵ
1.0	Nonrel.	147.62	-0.0050743	0.10303
1.0	Rel.	147.62	-0.0050744	0.10304
10.0	Nonrel.	102.71	-0.66593	1.1267
10.0	Rel.	102.71	-0.66593	1.1267
50.0	Nonrel.	62.929	-6.3189	2.0853
50.0	Rel.	62.929	-6.3189	2.0854
100.0	Nonrel.	43.531	-12.093	2.4899
100.0	Rel.	43.531	-12.093	2.4899
350.0	Nonrel.	2.6451	-26.719	4.9223
350.0	Rel.	2.6452	-26.719	4.9222

III. THE CONSTRUCTION OF THE NN FORCE EMBEDDED IN THE 3N SYSTEM

If one defines $2\omega_p(\hat{k}) \equiv \sqrt{4(\omega(\hat{k}))^2 + p^2}$ then using (4) Eq. (8) can be written

$$\begin{aligned} & (v_p + 2\omega_p(\hat{k}))^2 = (2\omega(\hat{k}) + v)^2 + p^2 \\ = \; & 4(\omega_p(\hat{k}))^2 + 4mV \end{aligned} \tag{22}$$

or

$$(v_p)^2 + 2v_p\omega_p(\hat{k}) + 2\omega_p(\hat{k})v_p = 4mV \tag{23}$$

Between free states it yields

$$\begin{aligned} & v_p(k,k\,')(\omega_p(k) + \omega_p(k\,')) \\ = \; & 2mV(k,k\,') \\ - \; & \frac{1}{2}\int_0^\infty dk\,''(k\,'')^2 v_p(k,k\,'')v_p(k\,'',k\,') \end{aligned} \tag{24}$$

This has the same structure as (14) replacing $\omega(k)$ by $\omega_p(k)$. The iteration procedure described in section II works equally well. We compare in Fig 4 $v(k,k\,')$ and $v_p(k,k\,')$. We see a weakening of v_p against v. This is a fact known from previous calculations [13] and from the approximate (but very useful) expression [7]

$$v_p(k,k\,') = v(k,k\,')(1 - \frac{p^2}{8\omega(k)\omega(k\,')}) \tag{25}$$

This is demonstrated in Fig 4, where also an even simpler approximation

$$v_p(k,k\,') = v(k,k\,')(1 - \frac{p^2}{8m^2}) \tag{26}$$

is shown. That latter approximate version, however, is somewhat worse and is not recommended.

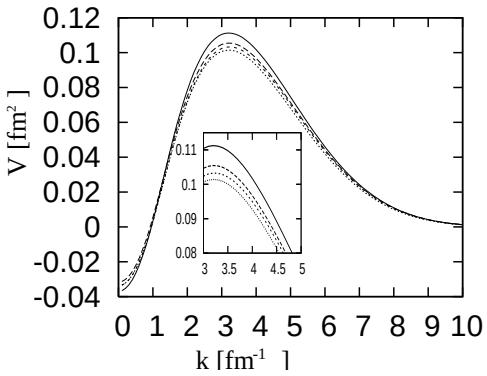

FIG. 4: Comparison $v(k,k\,')$ against $v_p(k,k\,')$ for AV18 in the state $^1S_0$ for fixed $k' = 1.0$ fm^{-1}. The solid curve shows the relativistic potential v ($p = 0$). The other curves show v_p for $p = 4.0$fm^{-1}. The long dashed curve is the exact one, the short dashed and dotted curves show the approximations for v_p given in Eqs.(25) and (26), respectively.

IV. SUMMARY

Relativistic calculations in the instant form of dynamics proposed in [15] requires an interacting 3N mass operator. As has been shown in [8] and used in [6],[7],[13], [16], [17], and [18] the NN potential in a moving frame v_p enters in the form given in Eq.(8), where the NN force in the NN c.m.system v enters into the relativistic NN Schrödinger equation (2). We showed that the quadratic operator relation (4) for v can be solved directly in an iterative manner and this very precisely. This has been documented by evaluating NN phase shift and mixing parameters using the standard nonrelatvistic Schrödinger equation (1) and the relativistic NN Schrödinger equation (2). This opens the way to get v's related to any NN potential V adjusted in a nonrelativistic frame work like the high precision NN potentials. By the same iterative procedure also v_p can be gained. It turned out that v is smaller in magnitude than V and v_p is smaller than v. Applications to the 3N bound and scattering states are planned.

The numerical calculations has been performed on the IBM Regatta p690+ of the NIC in Jülich, Germany.

V. APPENDIX

Derivation of Eq.(11)
We start from (9) and use the well known decomposition

$$< k|\Psi'_k >= \frac{\delta(k-k')}{kk\,'} + \frac{T(k,k';\frac{k'^2}{m})}{\frac{k'^2}{m} + i\epsilon - \frac{k^2}{m}} \tag{27}$$

to arrive at

$$< k|v|k\,' >= \Psi_b(k) M_b \Psi_b(k\,')$$

$$
\begin{aligned}
&+ \frac{T^*(k',k;\frac{k^2}{m})}{\frac{k^2}{m}-i\epsilon-\frac{k'^2}{m}}2\omega(k)\\
&+ \frac{T(k,k';\frac{k'^2}{m})}{\frac{k'^2}{m}+i\epsilon-\frac{k^2}{m}}2\omega(k')\\
&+ \int_0^\infty dk''(k'')^2\frac{T(k,k'';\frac{k''^2}{m})}{\frac{k''^2}{m}+i\epsilon-\frac{k^2}{m}}2\omega(k'')\\
&\frac{T^*(k',k'';\frac{k''^2}{m})}{\frac{k''^2}{m}-i\epsilon-\frac{k'^2}{m}} \qquad (28)
\end{aligned}
$$

The integral requires some care and we keep the limiting processes for the two scattering states separately by putting

$$
\begin{aligned}
&\frac{1}{k''^2-k^2+i\epsilon}\frac{1}{k''^2-k'^2-i\epsilon}\\
\to\ &\frac{1}{k''^2-k^2+i\epsilon_1}\frac{1}{k''^2-k'^2-i\epsilon_2}\\
=\ &\left(\frac{1}{k''^2-k^2+i\epsilon_1}-\frac{1}{k''^2-k'^2-i\epsilon_2}\right)\times\\
&\frac{1}{k^2-k'^2-i(\epsilon_1+\epsilon_2)}. \qquad (29)
\end{aligned}
$$

This allows us to perform one limit firstly with the result

$$
\begin{aligned}
&\frac{1}{k''^2-k^2+i\epsilon}\frac{1}{k''^2-k'^2-i\epsilon}\\
\to\ &\left(\frac{\mathcal{P}}{k''^2-k^2}-i\pi\delta(k''^2-k^2)\right)\frac{1}{k^2-k'^2-i\epsilon_2}\\
-\ &\left(\frac{\mathcal{P}}{k''^2-k'^2}+i\pi\delta(k''^2-k'^2)\right)\frac{1}{k^2-k'^2-i\epsilon_1} \qquad (30)
\end{aligned}
$$

Thus we get for some well behaved function $f(k'')$

$$
\begin{aligned}
&\int_0^\infty dk''(k'')^2\frac{f(k'')}{(k''^2-k^2+i\epsilon_1)(k''^2-k'^2-i\epsilon_2)}\\
&=\lim_{\epsilon\to+0}\frac{1}{k^2-k'^2-i\epsilon}\times\\
&\left(\mathcal{P}\int_0^\infty dk''(k'')^2\frac{f(k'')}{k''^2-k^2}-\mathcal{P}\int_0^\infty dk''(k'')^2\frac{f(k'')}{k''^2-k'^2}\right)\\
&-i\pi\lim_{\epsilon\to+0}\frac{1}{k^2-k'^2-i\epsilon}\times\\
&(\int_0^\infty dk''(k'')^2f(k'')\delta(k''^2-k^2)\\
+\ &\int_0^\infty dk''(k'')^2f(k'')\delta(k''^2-k'^2)) \qquad (31)
\end{aligned}
$$

The principal value prescription is denoted as "$\mathcal{P}\int$". In our case

$$
f(k'')=2m^2\omega(k'')T(k,k'';\frac{k''^2}{m})T^*(k',k'';\frac{k''^2}{m}) \qquad (32)
$$

and therefore

$$
\begin{aligned}
&\int_0^\infty dk''(k'')^2f(k'')\delta(k''^2-k^2)\\
=\ &2m^2\int_0^\infty dk''(k'')^2\omega(k'')T(k,k'';\frac{k''^2}{m})\\
&T^*(k',k'';\frac{k''^2}{m})\delta(k''^2-k^2)\\
=\ &m^2k\omega(k)T(k,k;\frac{k^2}{m})T^*(k',k;\frac{k^2}{m}) \qquad (33)
\end{aligned}
$$

This is part of the unitary relation

$$
\begin{aligned}
&T(k,k';\frac{q^2}{m})-T^*(k',k;\frac{q^2}{m})\\
=\ &2i\Im T(k,k';\frac{q^2}{m})\\
=\ &2i\Im T(k',k;\frac{q^2}{m})\\
-\ &2i\pi m\int_0^\infty dk''(k'')^2T(k,k'';\frac{q^2}{m})T^*(k',k'';\frac{q^2}{m})\delta(q^2-k''^2)\\
=\ &-i\pi mqT(k,q;\frac{q^2}{m})T^*(k',q;\frac{q^2}{m}) \qquad (34)
\end{aligned}
$$

(Note we used the symmetry of the off- the- energy- shell T-matrix). Consequently

$$
\begin{aligned}
&\int_0^\infty dk''k''^2f(k'')\delta(k''^2-k^2)\\
=\ &-\frac{2\omega(k)m}{\pi}\Im[T(k,k';\frac{k^2}{m})] \qquad (35)
\end{aligned}
$$

and

$$
\begin{aligned}
&-i\pi\lim_{\epsilon\to+0}\frac{1}{k^2-k'^2-i\epsilon}\times\\
&(\int_0^\infty dk''(k'')^2f(k'')\delta(k''^2-k^2)\\
+\ &\int_0^\infty dk''(k'')^2f(k'')\delta(k''^2-k'^2))\\
&=\lim_{\epsilon\to+0}\frac{2mi}{k^2-k'^2-i\epsilon}\times\\
&\left(\omega(k)\Im[T(k,k';\frac{k^2}{m})]+\omega(k')\Im[T(k',k;\frac{k'^2}{m})]\right) \qquad (36)
\end{aligned}
$$

Combined with Eq. (28) certain terms cancel and one arrives at Eq.(11).

[1] R. B. Wiringa, V. G. J. Stoks,and R. Schiavilla, Phys. Rev. C **51**,38 (1995).

[2] R. Machleidt, Phys. Rev. C **63**,024001 (2001).

[3] V. G. J. Stoks, R. A. M. Klomp, C. P. F. Terheggen, and

J. J. de Swart, Phys. Rev. C **49**, 2950 (1994).
[4] F. Coester, S.C. Pieper, F.J.D. Serduke, Phys. Rev C **11**, 1 (1975).
[5] J.L. Friar, Phys.Rev C **60**, 034002 (1999).
[6] W.Glöckle, T.-S.-H.Lee,F.Coester, Phys.Rev C **33**, 709 (1986).
[7] H. Witala, J. Golak, W. Glöckle,and H. Kamada, Phys.Rev. C **71**,054001 (2005).
[8] F.Coester, Helv.Phys. Acta **38**, 7 (1965).
[9] J.Carlson, V.R. Pandharipande,and R.Schiavilla,Phys.Rev C **47**, 484(1993).
[10] J.L.Forest, V.R.Pandharipande, J.Carlson, and R.Schiavilla, Phys.Rev C **52**, 576(1995).
[11] H. Kamada and W. Glöckle, Phys. Rev. Lett. **80**, 2547 (1998)
[12] T.W. Allen, G.L. Payne, W.N. Polyzou, Phys. Rev. C **62**,054002 (2000).
[13] H. Kamada, W. Glöckle, J. Golak, Ch. Elster, Phys. Rev. C **66**, 044010 (2002).
[14] R. Machleidt, Adv. Nucl. Phys. **19** (1989) 189.
[15] P. A. M. Dirac, Rev. Mod. Phys. **21** (1949) 392.
[16] T. Lin, Ch. Elster, W. N. Polyzou, W. Glöckle, Phys. Rev. C**76** (2007) 044010.
[17] H. Witała, J. Golak, R. Skibiński, Phys. Lett. **B 634** (2006).
[18] R. Skibiński, H. Witała, J. Golak, nucl-th/0604033